

# Lepton flavor violating decay of SM-like Higgs in a radiative neutrino mass model

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## Abstract

The lepton flavor violating decay of the Standard Model-like Higgs (LFVHD) is discussed in the framework of the radiative neutrino mass model built in [17]. The branching ratio (BR) of the LFVHD are shown to reach  $10^{-5}$  in the most interesting region of the parameter space shown in [17]. The dominant contributions come from the singly charged Higgs mediations, namely the coupling of  $h_2^\pm$  with exotic neutrinos. Furthermore, if doubly charged Higgs is heavy enough to allow the mass of  $h_2^\pm$  around 1 TeV, the mentioned BR can reach  $10^{-4}$ . Besides, we have obtained that the large values of the  $\text{Br}(h \rightarrow \mu\tau)$  leads to very small ones of the  $\text{Br}(h \rightarrow e\tau)$ , much smaller than various sensitivity of current experiments.

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## I. INTRODUCTION

The confirmation of the existence of a scalar boson, known as the Standard Model (SM)-like Higgs, is the greatest early success of the LHC [1, 2]. In addition, the LHC has reported recently some significant new physics beyond the SM where the LFVHD  $h \rightarrow \mu\tau$  is one of the hottest subjects [3]. The upper bound  $\text{Br}(h^0 \rightarrow \mu\tau) < 1.5 \times 10^{-2}$  at 95% C.L. was announced by the CMS Collaboration, in agreement with  $1.85 \times 10^{-2}$  at 95% C.L. from the ATLAS Collaboration. More interestingly, the CMS has indicated a  $2\sigma$  excess of this decay, with the value of  $\text{Br}(h \rightarrow \mu\tau) = 0.84^{+0.39}_{-0.37}\%$ . Besides, two other lepton flavor violating (LFV) decays of the SM-like Higgs have set experimental upper bounds at  $\text{BR}(h \rightarrow e\tau) < 0.7\%$  and  $\text{BR}(h \rightarrow e\mu) < 0.036\%$  at 95% C.L.[4]. Theoretically, many publications have studied how large the  $\text{BR}(h \rightarrow \mu\tau)$  can become in specific models beyond the SM, such as the seesaw [5, 6], supersymmetry (SUSY)[5, 7], two Higgs Doublet [8, 9], and 3-3-1 models [10], as well as other interesting ones [11–13]. The LFV decay of new neutral Higgs bosons in non-SUSY models has also been discussed [14]. The significance of LFVHD in colliders was addressed in [15].

The first source of LFV decays results from the mixing of different flavor massive neutrinos [16]. The simplest models explaining the mixing and masses of active neutrinos may be the seesaw models, but the BR of LFV decays predicted by these models are very small. Perhaps the inverse seesaw model gives the largest BR, which is about  $10^{-5}$  [6]. All of the SUSY models, even the Minimal Supersymmetric Standard Model, easily predict large values of the BR of LFVHD with new LFV sources in the slepton sector. However, the particle spectra of these models are rather complicated. In contrast, recent studies have shown that many of non-SUSY models inheriting simpler particle spectra can predict very large BRs of the LFVHD at one loop level, satisfying all relevant experimental constraints. Some of these models even have tree level couplings of LFVHD, and they simultaneously explain other interesting experimental results [8].

There is another class of models, where neutrino mass is radiatively generated, that can predict large BRs of the LFVHD. These models do not have active neutrino mass terms at tree level, but they contain LFV couplings of new particles such as scalars and new leptons in order to generate neutrino masses from loop contributions. There is an interesting property that the loop suppression factors appearing in the expression of neutrino masses lead to

the alleviation of the hierarchy in couplings. Hence the aforementioned models will allow large Yukawa couplings, which may result in large BR values for many LFV processes. By investigating a specific model with three loop neutrino mass introduced in [17], we try to make clear how large the BR of LFVHD can reach in the allowed regions. Furthermore, the contributions from active neutrino mediations may be enhanced because the Glashow-Iliopoulos-Maiani (GIM) mechanism does not work. Using the 't Hooft-Feynman gauge where loop contributions from private Feynman diagrams are all finite, we can compute and compare them. As a result, the best regions for large BRs of the LFVHD can be found with precise conditions of free parameters. The contributions from active neutrino mediations are divided completely into independent contributions of  $W^\pm$  and new singly charged Higgs bosons. As we will see later, the active neutrino loops in the radiative neutrino mass model may give significant contributions to LFV processes. This is different from all of the other models, where these contributions are either ignored; or are difficult to estimate when active neutrinos mix with new leptons, as in the case of the (inverse) seesaw models.

Our paper is arranged as follows. Section II will collect all needed ingredients for calculating the BR of the LFVHD. Section III concentrates on detailed expressions of the LFVHD amplitudes and partial widths. The constraints given in [17] will be discussed to find the allowed regions of parameter space. A numerical discussion is conducted and the main results are summarized in Secs. IV and V. Finally, the Appendices A and B list analytic expressions of Passarino-Veltman (PV) functions and LFVHD form factors. The divergence cancellations of particular one-loop Feynman diagrams in the 't Hooft-Feynman gauge are proved in Appendix B.

## II. REVIEW OF THE MODEL

### A. Particle content

Following Ref. [17], the particle content of the model is listed in Table I, where the last row represents charges of an additional global symmetry,  $U(1)$ . Aside from the SM particles, new particles are all gauge singlets, including three Majorana fermions,  $N_{R_{1,2,3}}$ ; one neutral Higgs boson  $\Sigma_0$ ; four singly charged Higgs bosons ( $h_1^\pm, h_2^\pm$ ); and two doubly charged Higgses,  $k^{\pm\pm}$ . After the breaking of the  $U(1)$  symmetry, a remnant  $Z_2$  symmetry keeps  $N_{R_{1,2,3}}$  and  $h_2^\pm$

as negative parity particles. The remaining particles are trivial. An interesting consequence is that the lightest Majorana neutrino, which has negative parity, will be stable and can be a dark matter candidate.

	Lepton Fields			Scalar Fields				
Characters	$L'_{L_i} = \begin{pmatrix} \nu'_{L_i} \\ e_{L_i} \end{pmatrix}$	$e_{R_i}$	$N_{R_i}$	$\Phi$	$\Sigma_0$	$h_1^+$	$h_2^+$	$k^{++}$
$SU(3)_C$	<b>1</b>	<b>1</b>	<b>1</b>	<b>1</b>	<b>1</b>	<b>1</b>	<b>1</b>	<b>1</b>
$SU(2)_L$	<b>2</b>	<b>1</b>	<b>1</b>	<b>2</b>	<b>1</b>	<b>1</b>	<b>1</b>	<b>1</b>
$U(1)_Y$	-1/2	-1	0	1/2	0	1	1	2
$U(1)$	0	0	$-x$	0	$2x$	0	$x$	$2x$

TABLE I: Lepton and scalar fields proposed in [17]. The notations  $L'_{L_i}$  and  $\nu'_{L_i}$  denote the flavor states in distinguish with the mass state  $\nu_{L_i}$  used later.

The Yukawa sector  $\mathcal{L}_Y$  respecting all mentioned symmetries is given as

$$-\mathcal{L}_Y = (y_\ell)_{ij} \overline{L'_{L_i}} \Phi e_{R_j} + \frac{1}{2} (y_L)_{ij} \overline{(L'_{L_i})^c} L'_{L_j} h_1^+ + (y_R)_{ij} \overline{N_{R_i}} (e_{R_j})^c h_2^- + \frac{1}{2} (y_N)_{ij} \Sigma_0 \overline{(N_{R_i})^c} N_{R_j} + \text{h.c.} \quad (1)$$

In addition, when symmetries are broken an effective Yukawa term appears after we take into account the loop contributions for generating active neutrino masses, namely,

$$-\mathcal{L}_Y^{\text{eff}} = \frac{(m_\nu)_{ab}}{\sqrt{2}v'} \times \overline{(\nu_{L_a})^c} \nu_{L_b} \times \Sigma_0, \quad (2)$$

corresponding to the active neutrino mass term derived in [17].

The Higgs potential is

$$\begin{aligned} \mathcal{V} = & m_\Phi^2 |\Phi|^2 + m_\Sigma^2 |\Sigma_0|^2 + m_{h_1}^2 |h_1^+|^2 + m_{h_2}^2 |h_2^+|^2 + m_k^2 |k^{++}|^2 \\ & + \left[ \lambda_{11} \Sigma_0^* h_1^- h_1^- k^{++} + \mu_{22} h_2^+ h_2^+ k^{--} + \text{h.c.} \right] + \lambda_\Phi |\Phi|^4 + \lambda_{\Phi\Sigma} |\Phi|^2 |\Sigma_0|^2 + \lambda_{\Phi h_1} |\Phi|^2 |h_1^+|^2 \\ & + \lambda_{\Phi h_2} |\Phi|^2 |h_2^+|^2 + \lambda_{\Phi k} |\Phi|^2 |k^{++}|^2 + \lambda_\Sigma |\Sigma_0|^4 + \lambda_{\Sigma h_1} |\Sigma_0|^2 |h_1^+|^2 + \lambda_{\Sigma h_2} |\Sigma_0|^2 |h_2^+|^2 \\ & + \lambda_{\Sigma k} |\Sigma_0|^2 |k^{++}|^2 + \lambda_{h_1} |h_1^+|^4 + \lambda_{h_1 h_2} |h_1^+|^2 |h_2^+|^2 + \lambda_{h_1 k} |h_1^+|^2 |k^{++}|^2 \\ & + \lambda_{h_2} |h_2^+|^4 + \lambda_{h_2 k} |h_2^+|^2 |k^{++}|^2 + \lambda_k |k^{++}|^4. \end{aligned} \quad (3)$$

The scalar fields are parameterized as

$$\Phi = \begin{pmatrix} G_w^+ \\ \frac{v+\phi+iG_z}{\sqrt{2}} \end{pmatrix}, \quad \Sigma_0 = \frac{v' + \sigma}{\sqrt{2}} e^{iG/v'}, \quad (4)$$

where  $v \simeq 246$  GeV,  $v'$  is a new vacuum expectation value (VEV), and  $G_w^\pm$  and  $G_z$  are Goldstone bosons of  $W^\pm$  and  $Z$  bosons, respectively.

The above lepton and Higgs sectors show that the singly charged Higgs bosons contribute mainly to the BR of the LFBVD while the doubly charged Higgs bosons do not. Because all of charged Higgs bosons are  $SU(2)_L$  singlets, they do not couple with  $W$  bosons like new Higgs multiplets in other models. At the one loop level and t' Hooft-Feynman gauge, the Feynman diagrams for LFBVD are shown in Fig. 1.

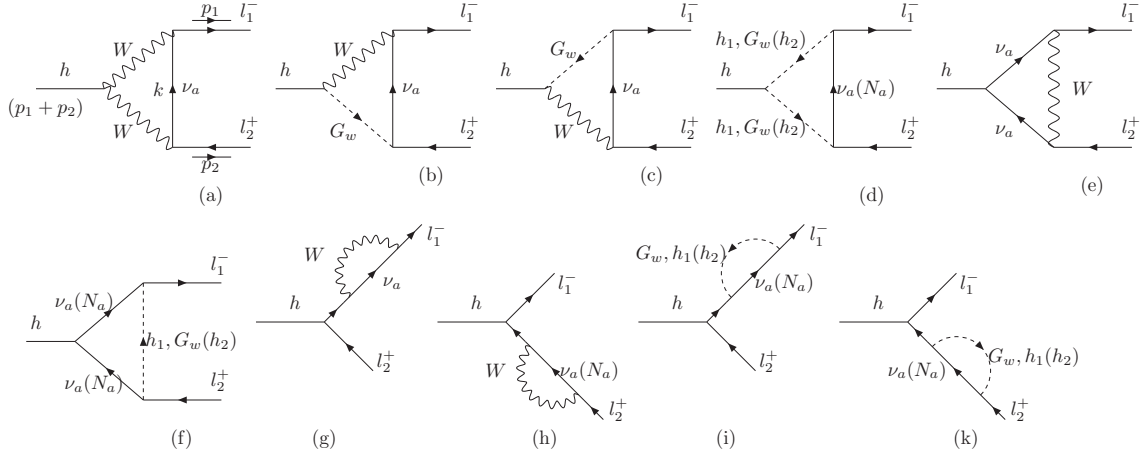


FIG. 1: Feynman diagrams for LFBVD  $h \rightarrow \mu^\pm \tau^\mp$  decay, where  $h, l_1 \equiv e_2 = \mu, l_2 \equiv e_3 = \tau$ . The parentheses imply that  $N_a$  and  $h_2^\pm$  couple with each other only.

In order to investigate the LFBVD, the amplitudes will be calculated using the t' Hooft-Feynman gauge. The amplitudes are formulated as functions of the PV functions analyzed in [10]. The notations of four-component (Dirac) spinors are used for leptons. Specifically, the charged leptons are  $e_i$ , so that the left-handed component is  $e_{L_i} = P_L e_i$  and right-handed one is  $e_{R_i} = P_R e_i$ , where  $P_{R,L} \equiv (1 \pm \gamma_5)/2$  are chiral operators. The corresponding charge conjugations are  $e_i^c \equiv C \bar{e}_i^T = (e_{L_i}^c, e_{R_i}^c)^T$ , which satisfies  $e_{L_i}^c = P_L e_i^c = C \bar{e}_{R_i}^T \equiv (e_{R_i})^c$  and  $e_{R_i}^c = P_R e_i^c = C \bar{e}_{L_i}^T \equiv (e_{L_i})^c$ . For the Majorana leptons such as active and exotic neutrinos, the four-spinors are  $\nu_i' = \nu_i'^c = (\nu_{L_i}', (\nu_{L_i}')^c)^T = (\nu_{L_i}', \nu_{R_i}'^c)^T$  and  $N_i = N_i^c = ((N_{R_i})^c, N_{R_i})^T = (N_{L_i}^c, N_{R_i})^T$ . To reduce the second and third terms of (1) to more

convenient forms, we will use equalities like  $\overline{e_i^c} P_L \nu_j' = \overline{\nu_j^c} P_L e_i = \overline{\nu_j'} P_L e_i$  and  $\overline{N_{R_i}}(e_{R_j})^c = \overline{N_i} P_L e_j^c = \overline{e_j} P_L N_i^c = \overline{e_j} P_L N_i$  (Appendix G of [20]).

## B. Mass spectrum and LFVHD couplings

In the mass basis, the model consists of two CP-even neutral Higgs bosons; a Nambu-Goldstone scalar boson  $G$ ; four singly charged Higgs bosons,  $h_1^\pm, h_2^\pm$ ; and two doubly charged Higgs bosons,  $k^{\pm\pm}$ . At one loop level, the LFVHD involves only two CP-even neutral and four singly charged Higgs bosons. The two CP-even neutral Higgs bosons are a SM-like Higgs,  $h$  ( $m_h = 125$  GeV) and a new CP-even one,  $H$ . They relate to the original Higgs components  $(\phi, \sigma)$  via the transformation

$$\begin{pmatrix} \phi \\ \sigma \end{pmatrix} = \begin{pmatrix} c_\alpha & s_\alpha \\ -s_\alpha & c_\alpha \end{pmatrix} \begin{pmatrix} h \\ H \end{pmatrix}, \quad (5)$$

where  $c_\alpha \equiv \cos \alpha$ ,  $s_\alpha \equiv \sin \alpha$ , and  $\alpha$  is defined as

$$\sin 2\alpha = \frac{2\lambda_{\Phi\Sigma} v v'}{m_H^2 - m_h^2}. \quad (6)$$

The masses  $m_h$  and  $m_H$  are functions of the four parameters  $\lambda_\Phi, \lambda_\Sigma, \lambda_{\Phi\Sigma}$  and  $v'$ . In the following calculations, we will fix  $m_h = 125.1$  GeV while  $m_H, s_\alpha$  and  $v'$  are taken as free parameters. The original parameters are given as [17]

$$\lambda_{\Phi\Sigma} = \frac{s_\alpha c_\alpha (m_H^2 - m_h^2)}{v v'}, \quad \lambda_\Phi = \frac{c_\alpha^2 m_h^2 + s_\alpha m_H^2}{2v^2}, \quad \text{and} \quad \lambda_\Sigma = \frac{s_\alpha^2 m_h^2 + c_\alpha^2 m_H^2}{2v'^2}. \quad (7)$$

The perturbativity limit forces  $\lambda_\Phi, \lambda_\Sigma, \lambda_{\Phi\Sigma} \leq 4\pi$ , therefore gives an upper bound on  $m_H$  with a large  $s_\alpha$ . For example,  $|s_\alpha| \leq 0.3$  corresponds to  $m_H \leq 4$  TeV.

The masses of the singly charged Higgs bosons are given as

$$m_{h_1^\pm}^2 = m_{h_1}^2 + \frac{1}{2} (\lambda_{\Phi h_1} v^2 + \lambda_{\Sigma h_1} v'^2), \quad m_{h_2^\pm}^2 = m_{h_2}^2 + \frac{1}{2} (\lambda_{\Phi h_2} v^2 + \lambda_{\Sigma h_2} v'^2). \quad (8)$$

Regarding the lepton sector, the first and last terms in (1) correspond to the mass terms of charged leptons and exotic neutrinos, respectively. They are assumed to be diagonal, i.e., the flavor basis and mass basis coincide. Their expressions are obtained as

$$m_{e_i} = (y_\ell)_{ii} \frac{v}{\sqrt{2}}, \quad m_{N_i} = \frac{v'}{\sqrt{2}} (y_N)_{ii}. \quad (9)$$

The active neutrino masses generated from three loop corrections can be expressed by the effective term (2). The components of active neutrino mass matrix are given as [17]

$$(m_\nu)_{ab} = \frac{\mu_{11}\mu_{22}}{(4\pi)^6} \sum_{i,j,k=1}^3 \frac{1}{M_k^4} [(y_L)_{ai} m_{e_i} (y_R^T)_{ik} (M_{N_k}) (y_R)_{kj} m_{e_j} (y_L^T)_{jb}] \\ \times F_1 \left( \frac{m_{h_1^\pm}^2}{M_k^2}, \frac{m_{h_2^\pm}^2}{M_k^2}, \frac{m_{\ell_i}^2}{M_k^2}, \frac{m_{\ell_j}^2}{M_k^2}, \frac{M_{N_k}^2}{M_k^2}, \frac{m_{k^{\pm\pm}}^2}{M_k^2} \right), \quad (10)$$

where  $M_k$  is the maximal value among the quantities  $m_{h_1^\pm}, m_{h_2^\pm}, m_{e_i}, m_{e_j}, M_{N_k}, m_{k^{\pm\pm}}$ ;  $\mu_{11} \equiv \lambda_{11} v' / \sqrt{2}$ ; and  $F_1$  is the three loop function given in detail in [17].

All of the mass terms- together with couplings  $h\bar{f}f$ , where  $f = e_i, \nu_i$  and  $N_i$ - are parts of Yukawa terms of neutral Higgs bosons. The flavor states  $\nu'_{L_i}$  and mass states  $\nu_{L_i}$  ( $i = 1, 2, 3$ ) of active neutrinos are related by the transformation  $\nu'_{L_i} = U_{ij}^L \nu_{L_j}$ , where  $U^L U^{L\dagger} = U^{L\dagger} U^L = 1$ . The masses and mixing angles of the active neutrinos are taken from the best-fit experimental data given in [19]. The only unknown parameter is the lightest mass.

Concerning only on the mass terms and couplings involving the LFBVHD of the SM-like Higgs boson, the Yukawa interactions (1) and (2) are written as follows:

$$- \mathcal{L}'_Y = m_{e_i} \bar{e}_i e_i + \frac{1}{2} m_{\nu_i} \bar{\nu}_i \nu_i + \frac{1}{2} m_{N_i} \bar{N}_i N_i + \frac{m_{e_i}}{v} \bar{e}_i e_i (c_\alpha h) - \left[ \frac{m_{N_i}}{2v'} \bar{N}_i N_i + \frac{m_{\nu_i}}{2v'} \bar{\nu}_i \nu_i \right] s_\alpha h \\ + \frac{\sqrt{2} m_{e_i}}{v} [U_{ij}^L \bar{e}_i P_L \nu_j G_w^- + U_{ij}^{L*} \bar{\nu}_j P_R e_i G_w^+] + (y_L^T U^{L*})_{ij} \bar{\nu}_j P_L e_i h_1^+ + (y_L^T U^L)_{ij} \bar{e}_i P_R \nu_j h_1^- \\ + (y_R^T)_{ij} \bar{e}_i P_L N_j h_2^- + (y_R^T)_{ij} \bar{N}_j P_R e_i h_2^+. \quad (11)$$

The LFBV couplings relating to the  $W^\pm$  gauge boson only occur in the covariant kinetic terms of  $SU(2)_L$  doublets, exactly the same as in the SM, namely,

$$\mathcal{L}_{\text{kin}}^l = i \bar{L}_{L_i} \gamma^\mu D_\mu L_{L_i} + (D_\mu \Phi)^\dagger (D^\mu \Phi), \quad (12)$$

where  $D_\mu$  is the covariant derivative defined in the SM. All relevant couplings of the LFBVHD are collected in Table II.

### C. Parameter constraints from the previous work

For calculating the BR of the LFBVHD, in the following sections we will mainly use the constraints of parameters obtained in [17]. The important points are reviewed as follows. The parameters in the model were first investigated to ensure that they satisfy the neutrino

Vertex	Coupling	Vertex	Coupling
$h\bar{e}_i e_i$	$-\frac{im_{e_i}}{v}c_\alpha$	$h\bar{\nu}_i \nu_i$	$\frac{im_{\nu_i}}{v'}s_\alpha$
$h\bar{N}_i N_i$	$\frac{im_{N_i}}{v'}s_\alpha$	$hW_\mu^+ W_\nu^-$	$igm_W c_\alpha g^{\mu\nu}$
$hh_1^+ h_1^-$	$i(-vc_\alpha \lambda_{\Phi h_1} + v' s_\alpha \lambda_{\Sigma h_1})$	$hh_2^+ h_2^-$	$i(-vc_\alpha \lambda_{\Phi h_2} + v' s_\alpha \lambda_{\Sigma h_2})$
$h(p_0)W_\mu^+ G_w^-(p_-)$	$\frac{ig}{2}c_\alpha(p_0 - p_-)^\mu$	$h(p_0)W_\mu^- G_w^+(p_+)$	$-\frac{ig}{2}c_\alpha(p_0 - p_+)^\mu$
$\bar{e}_i \nu_j h_1^-$	$-i(y_L^T U^L)_{ij} P_R$	$\bar{\nu}_j e_i h_1^+$	$i(y_L^T U^L)_{ij}^* P_L$
$\bar{e}_i \nu_j G_w^-$	$-i\frac{\sqrt{2}m_{e_i}}{v}U_{ij}^L P_L$	$\bar{\nu}_j e_i G_w^+$	$-i\frac{\sqrt{2}m_{e_i}}{v}U_{ij}^{L*} P_R$
$\bar{e}_i N_j h_2^-$	$-i(y_R^T)_{ij} P_L$	$\bar{N}_j e_i h_2^+$	$-i(y_R^T)_{ij} P_R$
$\bar{e}_i \nu_j W_\mu^-$	$\frac{ig}{\sqrt{2}}U_{ij}^L \gamma^\mu P_L$	$\bar{\nu}_j e_i W_\mu^+$	$\frac{ig}{\sqrt{2}}U_{ij}^{L*} \gamma^\mu P_L$
$hG_w^+ G_w^-$	$i(-2vc_\alpha \lambda_\Phi + v' s_\alpha \lambda_{\Phi\Sigma})$		

TABLE II: Couplings of LFVHD in the t' Hooft-Feynman gauge. The momenta are incoming

oscillation data, the current bounds of the BR of the LFV processes, the universality of charged currents, and the vacuum stability of the Higgs self-couplings. In addition, the doubly charged Higgses  $k^{\pm\pm}$  are assumed to be light enough that they could be detected at the LHC. The constraints on parameters involving with the LFVHD are (i) the Dirac phase of the active neutrino mixing matrix prefers the value of  $\delta = \pi$ , while the Majorana phase is still free; (ii) the masses of singly charged Higgs bosons should not be smaller than 3 TeV; (iii) the value of  $|(y_R)_{22}|$  should be around 9; (iv) the value of  $v'$  should be on the order of  $O(1)$  TeV. The investigation in [17] also showed that the heavier doubly charged Higgs bosons  $k^{\pm\pm}$  will allow the lighter singly charged Higgs bosons. This leads to an interesting consequence of large values of the BR of the LFVHD as we will show in the numerical investigation.

The constraint from the LHC Higgs boson search was also discussed in [17], including the effects of the  $U(1)$  global Goldstone boson in the invisible decay of the SM-like Higgs bosons and the pair annihilation of the dark matter (DM) candidate  $N_{R_1}$ . From this, the constraint of mixing angle of neutral Higgs bosons is obtained as  $|\sin\alpha| \leq 0.3$ . Finally, the condition of DM candidate mentioned above leads to the conclusion that the  $N_{R_1}$  mass should be around the value of  $m_h/2$  in order to successfully explain the current relic density of DM; the VEV  $v'$  was found smaller than 10 TeV.

Many other issues involving with global  $U(1)$  Goldstone boson were discussed in detail in



[17], for instance, anomaly-induced interaction to two photons, active-sterile neutral lepton mixing, and neutrinoless double beta decay via  $W$  exchange. Possible bounds from cosmological issue, such as the effect on cosmic microwave background via cosmic string generated by the spontaneous breaking down of the global  $U(1)$  symmetry, were also mentioned. None of these issues change the constraints of parameters indicated above. Though new constraints on DM masses in the presence of  $U(1)$  global symmetry were addressed in [29], more studies are needed for confirmation. Furthermore, the considered global symmetry can be moved straight to the local one [17], or replaced with a suitable discrete symmetry. This discussion is beyond the scope of this work.

In the next section, we will focus on parameters affecting the LFVHD and will discuss more clearly the relevant constraints if ones are needed.

### III. FORMULAS OF LFVHD AND PARAMETER CONSTRAINTS

The effective Lagrangian of the decay  $h \rightarrow \tau^\pm \mu^\mp$  is written as  $\mathcal{L}^{LFV} = h(\Delta_L \bar{\mu} P_L \tau + \Delta_R \bar{\mu} P_R \tau) + \text{h.c.}$ , where  $\Delta_{L,R}$  are scalar factors arising from the loop contributions. The decay amplitude is defined as  $i\mathcal{M} = i\bar{u}_1(\Delta_L P_L + \Delta_R P_R)v_2$  [6], where  $u_1 \equiv u_1(p_1, s_1)$  and  $v_2 \equiv v_2(p_2, s_2)$  are the Dirac spinors of a muon and a tauon, respectively. The partial width of the decay is given as

$$\Gamma(h \rightarrow \mu\tau) \equiv \Gamma(h \rightarrow \mu^- \tau^+) + \Gamma(h \rightarrow \mu^+ \tau^-) = \frac{m_h}{8\pi} (|\Delta_L|^2 + |\Delta_R|^2), \quad (13)$$

with the condition  $m_h \gg m_1, m_2$ , where  $m_1, m_2$  are muon and tauon masses, respectively. The on-shell conditions for external particles are  $p_i^2 = m_i^2$  (i=1,2) and  $p_h^2 \equiv (p_1 + p_2)^2 = m_h^2$ .

The loop contributions can be separated into two parts  $\Delta_L = \Delta_L^\nu + \Delta_L^N$  and  $\Delta_R = \Delta_R^\nu + \Delta_R^N$ , corresponding to the appearance of the active and exotic neutrinos in the loops. In the 't Hooft-Feynman gauge, the specific formulas of contributions from diagrams shown in Fig. 1 are listed in Appendix B, where new notations such as  $E_{L,R}$  factors are used. The

contribution of the loops with active neutrinos  $\Delta_L^\nu$  is obtained as

$$\begin{aligned}\Delta_L^\nu &= \frac{1}{16\pi^2} \sum_{a=1}^3 (y_L^T U^L)_{2a} (y_L^T U^L)_{3a}^* \left[ (v' \lambda_{hh_1 h_1}) E_L^{\nu h_1 h_1} + s_\alpha E_L^{h_1 \nu \nu} + (-c_\alpha) E_L^{h_1 \nu} \right] \\ &+ \frac{1}{16\pi^2} \sum_{a=1}^3 U_{2a}^L U_{3a}^{L*} \left[ g^3 c_\alpha E_L^{\nu WW} + \frac{g^2 c_\alpha}{2} (E_L^{\nu G_w W} + E_L^{\nu W G_w}) \right. \\ &\left. + (v' \lambda_{h G_w G_w}) E_L^{\nu G_w G_w} + s_\alpha \left( \frac{g^2}{2} E_L^{W \nu \nu} + E_L^{G_w \nu \nu} \right) - c_\alpha \left( \frac{g^2}{2} E_L^{W \nu} + E_L^{G_w \nu} \right) \right], \quad (14)\end{aligned}$$

where  $(v' \lambda_{hh_1 h_1}) = -v c_\alpha \lambda_{\Phi h_1} + v' s_\alpha \lambda_{\Sigma h_1}$  and  $(v' \lambda_{h G_w G_w}) = -2v c_\alpha \lambda_\Phi + v' s_\alpha \lambda_{\Phi \Sigma}$ .

The contribution from exotic neutrino mediations  $\Delta_L^N$  is given as

$$\Delta_L^N = \frac{1}{16\pi^2} \sum_{a=1}^3 (y_R^T)_{2a} (y_R^T)_{3a} \left[ (v' \lambda_{hh_2 h_2}) E_L^{N h_2 h_2} + s_\alpha E_L^{h_2 N N} - c_\alpha E_L^{N h_2} \right], \quad (15)$$

where  $(v' \lambda_{hh_2 h_2}) = -v c_\alpha \lambda_{\Phi h_2} + v' s_\alpha \lambda_{\Sigma h_2}$ .

Similarly, we have  $\Delta_R^\nu = \Delta_L^\nu (E_L \rightarrow E_R)$  and  $\Delta_R^N = \Delta_L^N (E_L \rightarrow E_R)$ .

As proved in Appendix B, the  $\Delta_{L,R}$  are convergent. Specifically, in the 't Hooft-Feynman gauge, the private contributions from specific diagrams are always finite. In addition, the limit  $p_1^2, p_2^2 \simeq 0$  result in the extremely small contributions of diagrams relating with the two point functions, namely 1g), 1h), 1i) and 1k). Hence their contributions are ignored. Besides, it can be estimated that the sum of contributions in the two last lines in (14) is very suppressed, because of the GIM mechanism, controlled by the factor  $\sum_{a=1}^3 U_{2a}^L U_{3a}^{L*}$ . This is a general property of all models where the neutrino masses are generated from the seesaw mechanism. Whereas the contributions in the first lines of (14) and (15) may be large because the appearance of the  $y_L$  and  $y_R$  breaks the GIM mechanism, non-zero contributions survive which do not contain factors of very light of neutrino masses.

In the numerical calculation, the following parameters are taken from experimental data, for example [22]:  $v \simeq 246$  GeV,  $m_h = 125.1$  GeV,  $m_W = 80.4$  GeV, and the muon and tauon masses are  $m_\mu = 0.105$  GeV,  $m_\tau = 1.776$  GeV. The total decay width of the SM-like Higgs bosons  $\Gamma_h = 4.1 \times 10^{-3}$  GeV is used. Based on the investigation of [17], relevant parameters of the active neutrino masses are only considered in the normal hierarchy scheme. In particular the mixing parameters  $U^L$  are expressed as follows

$$U^L = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{23} & s_{23} \\ 0 & -s_{23} & c_{23} \end{pmatrix} \begin{pmatrix} c_{13} & 0 & -s_{13} \\ 0 & 1 & 0 \\ s_{13} & 0 & c_{13} \end{pmatrix} \begin{pmatrix} c_{12} & s_{12} & 0 \\ -s_{12} & c_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (16)$$

where  $c_{ij} \equiv \cos \theta_{ij}$ ,  $s_{ij} \equiv \sin \theta_{ij}$  and the Dirac CP phase  $\delta$  and the Majorana CP phase  $\phi$  are taken as  $\delta = \pi$  and  $\phi = 0$ . The best-fit values of neutrino oscillation parameters given in [19]

$$\begin{aligned}\Delta m_{21}^2 &= 7.50 \times 10^{-5} \text{ eV}^2, & \Delta m_{31}^2 &= 2.457 \times 10^{-3} \text{ eV}^2, \\ s_{12}^2 &= 0.304, & s_{23}^2 &= 0.452, & s_{13}^2 &= 0.0218\end{aligned}\tag{17}$$

are also used in our numerical calculations. The lightest neutrino mass will be chosen as  $m_{\nu_1} = 10^{-10} \text{ GeV}$ , satisfying the condition  $\sum_b m_{\nu_b} \leq 0.5 \text{ eV}$  obtained from the cosmological constraint. The remain two neutrino masses are  $m_{\nu_b}^2 = m_{\nu_1}^2 + \Delta m_{\nu_b 1}^2$  where  $b = 2, 3$ .

The other unknown parameters involving the LFBVHD are: the VEV  $v'$ ; the mixing angle of the two neutral Higgs bosons  $\alpha$ ; the exotic neutrino masses  $m_{N_a}$ ; the new Higgs masses  $m_{h_1^\pm}, m_{h_2^\pm}, m_H$ ; the Yukawa coupling matrices  $y_L$  and  $y_R$ ; and the trilinear Higgs self-couplings  $\lambda_{\Phi h_1}, \lambda_{\Phi h_2}, \lambda_{\Sigma h_1}, \lambda_{\Sigma h_2}$ .

In the numerical investigation, we focus first on the most interesting regions of the parameter space indicated in [17], where  $m_{N_1}$  plays the role of a DM particle and doubly charged Higgses  $k^{\pm\pm}$  are light enough to be observed at the LHC. The values of the relevant parameters are summarized as follows:  $m_h/2 = m_{N_1} < m_{N_2} < m_{N_3}$ ,  $|s_\alpha| \leq 0.3$ ;  $3 \text{ TeV} \leq (m_{h_1^\pm}, m_{h_2^\pm}) \sim \mathcal{O}(1) \text{ TeV}$  and  $v' \sim \mathcal{O}(1) \text{ TeV}$ . The Yukawa coupling matrix  $y_L$  are satisfied the following conditions [23],

$$\begin{aligned}(y_L)_{13} &= \left( \frac{s_{12}c_{23}}{c_{12}c_{13}} - \frac{s_{13}s_{23}}{c_{13}} \right) (y_L)_{23} = 0.394(y_L)_{23}, \\ (y_L)_{12} &= \left( \frac{s_{12}s_{23}}{c_{12}c_{13}} + \frac{s_{13}c_{23}}{c_{13}} \right) (y_L)_{23} = 0.56(y_L)_{23},\end{aligned}\tag{18}$$

and  $|(y_L)_{23}| \leq 1$ . The conditions of gauge coupling universalities also imply that

$$|(y_L)_{23}| < \frac{0.3m_{h_1^\pm}}{1\text{TeV}}.\tag{19}$$

The most stringent constraint comes from the LFV decay of muon with  $\text{Br}(\mu \rightarrow e\gamma) < 5.7 \times 10^{-13}$ . It gives the direct upper bounds on the following products of the Yukawa couplings: (i)  $(y_L^T)_{23}(y_L^T)_{13}$  in the loop including virtual active neutrinos and  $h_1^\pm$  since  $y_L$  is antisymmetric; (ii)  $(y_R^T)_{2i}(y_R^T)_{1i}$  ( $i = 1, 2, 3$ ) in the loops including exotic neutrinos and  $h_2^\pm$ . The other constraints from the tauon decays are less stringent, hence are omitted here. On the other hand, the  $\text{Br}(h \rightarrow \mu\tau)$  depends on the products  $(y_R^T)_{2i}(y_R^T)_{3i}$  with

$i = 1, 2, 3$ . So, if  $|(y_R^T)_{ij}|$  with  $(ij) = \{(11), (21), (31)\}$  is small enough, the values of  $(y_R^T)_{ij}$  with  $(ij) = \{(22), (32), (33)\}$  may be large, without any inconsistency in the upper bounds of the BR in the LFV decays of charged leptons. In order to find the reasonable regions of parameter space, the upper bounds must be checked in the formula given in [17]:

$$\frac{\sum_{a=1}^3 \left[ (y_L^\dagger)_{af} (y_L)_{ia} \right]^2 (I_{1,a} I_{2,a} + I_{1,a}^2) + \sum_{a=1}^3 \left[ (y_R)_{af} (y_R^\dagger)_{ia} \right]^2 (I'_{1,a} I'_{2,a} + I_{1,a}^{\prime 2})}{16m_{h_1^\pm}^4 \left| \sum_{a=1}^3 (I_{1,a} I_{2,a} + I_{1,a}^2) \right|} < \frac{C_{if}}{[\text{TeV}]^4}, \quad (20)$$

where  $(i, f) = (\mu, e) = (2, 1)$ ,  $C_{if} = 1.6 \times 10^{-6}$  and many notations including  $(I_{1,a}, I_{2,a}, I'_{1,a}, I'_{2,a})$  are defined in detail in [17]. For simplicity, we mention only the following special cases. In the limit  $0 \simeq m_{\nu_a} \ll m_{h_1^\pm}$  and  $m_{h_2^\pm} \gg m_{N_1}$  we have

$$I_{1,a} \simeq -\frac{1}{(4\pi)^2} \frac{1}{36m_{h_1^\pm}^2}, \quad I_{2,a} \simeq -\frac{1}{(4\pi)^2} \frac{5}{36m_{h_1^\pm}^2}, \quad I'_{1,a} \simeq -\frac{1}{(4\pi)^2} \frac{1}{36m_{h_2^\pm}^2}, \quad I'_{2,a} \simeq -\frac{1}{(4\pi)^2} \frac{5}{36m_{h_2^\pm}^2}. \quad (21)$$

While a very large  $m_{N_a}$  gives  $I'_{1,a}, I'_{2,a} \simeq 0$ -for example when  $m_{N_{2,3}}$  are larger than few TeV-the bound (20) affects only the products  $|(y_L^\dagger)_{31} (y_L)_{23}|$  and  $|(y_R)_{31} (y_R^\dagger)_{23}|$ . Combining this with (18) we get new constraints:

$$|(y_L)_{23}| \leq 0.149 \frac{m_{h_1^\pm}}{1\text{TeV}}, \quad \text{and} \quad |(y_R)_{31} (y_R^\dagger)_{23}| < \left( \frac{0.1m_{h_2^\pm}}{1\text{TeV}} \right)^2. \quad (22)$$

This constraint of  $|(y_L)_{23}|$  is consistent with the numerical investigation done in [17], where  $m_{h_1^\pm} = 4.8 \text{ TeV}$  prefers  $(y_L)_{23} = 0.5 - 0.6 < 0.7 \simeq 4.8 \times 0.15$ . Interestingly, the first inequality in (22) is more strict than the one given in (19). The second constraint in (22) suggests that the small  $|(y_R)_{31}|$  will allow large  $|(y_R^\dagger)_{23}|$ , i.e the choice  $|(y_R)_{31}| \leq 10^{-3} \times (m_{h_2^\pm}/1\text{TeV})^2$  will allow  $|(y_R^\dagger)_{23}| \sim \mathcal{O}(1)$ . The  $m_{h_2^\pm} = 5(3) \text{ TeV}$  will give  $|(y_R)_{31}| \leq 0.25(0.09)$ , being consistent with the promoting regions indicated in [17]. But the absolute values of  $(y_R)_{22}$  and  $(y_R)_{23}$  should be smaller than the perturbative upper bound  $4\pi$ .

Another relevant constraint is the small upper bound of the  $\text{BR}(\mu \rightarrow 3e)$ . Following [17], all form factors relating with  $(y_R)_{ij}$  contain at least three factors  $(y_R)_{i1}$  or  $(y_R)_{1i}$ . Therefore they result in very suppressed values of the BR of the LFVHD if all  $|(y_R)_{i1,1i}|$ 's are small enough, without any conditions of small  $|(y_R)_{23,32,33}|$ 's.

Combining the above discussion with the analysis in [17], the reasonable values of the free parameters can be chosen as follows:  $s_\alpha = 0.3, \lambda \equiv 8\lambda_{\Phi h_1} = 8\lambda_{\Phi h_2} = \lambda_{\Sigma h_1} = \lambda_{\Sigma h_2} = 4$ ,  $m_{h_1^\pm} = m_{h_2^\pm} = 3 \text{ TeV}$ ,  $m_H = 2 \text{ TeV}$ ,  $(y_L)_{23} = \frac{0.14m_{h_1^\pm}}{1\text{TeV}}$ ,  $y'_R = (y_R)_{23} = (y_R)_{32} = (y_R)_{33} =$

3,  $(y_R)_{22} = 8$ ,  $v' \in \{1, 2, 8, 10\}$  TeV,  $m_{N_1} = m_h/2$ ,  $m_{N_2} = 1/2 m_{N_3} = 5$  TeV and  $(y_R)_{i,j} = 10^{-2}$  with at least one of the indices,  $i$  or  $j$  being 1. We would like to stress that the above choices are also based on the following additional reasons. The values of  $(y_L)_{23}$  and  $(y_R)_{ij}$  always satisfy all recent bounds of the BR of the LFV decays of charged leptons (22) as well as gauge coupling universalities (19). The  $\lambda$  parameter is positive and small enough to satisfy both conditions of perturbative limit and vacuum stability, whereas it is large enough to enhance the BR of LFVHD.

To investigate the variance of  $\text{BR}(h \rightarrow \mu\tau)$  versus the changing of free parameters, the ranges of free parameters will be chosen as follows:  $|s_\alpha| \leq 0.4$ ,  $0.1 < \lambda < 10$ ,  $|y'_R| \leq 4$ ,  $|(y_L)_{23}| \leq \frac{0.15 m_{h_1^\pm}}{1\text{TeV}}$ ,  $1\text{TeV} \leq m_{h_1^\pm}, m_{h_2^\pm}, v' < 10$  TeV and  $0.5\text{TeV} \leq m_{N_2} \leq 6$  TeV.

#### IV. NUMERICAL RESULT AND DISCUSSION

In this section, we will first investigate some private contributions to the BR of the LFVHD, namely the active neutrino loops with  $W^\pm/h_1^\pm$  bosons, and the exotic lepton loops. Based on this, the parameter space regions which give the large total contribution will be further studied.

The contributions of active neutrinos are shown in the Fig. 2. The left panel shows two contributions to the  $\text{BR}(h \rightarrow \mu\tau)$ : sum of all diagrams relating with the  $W^\pm$  gauge bosons and their Goldstone bosons, and Fig. 1f). The right panel shows the contribution to  $\text{BR}(h \rightarrow \mu\tau)$  from Fig. 1d) with virtual  $h_1^\pm$  in the loop. The figure emphasizes two

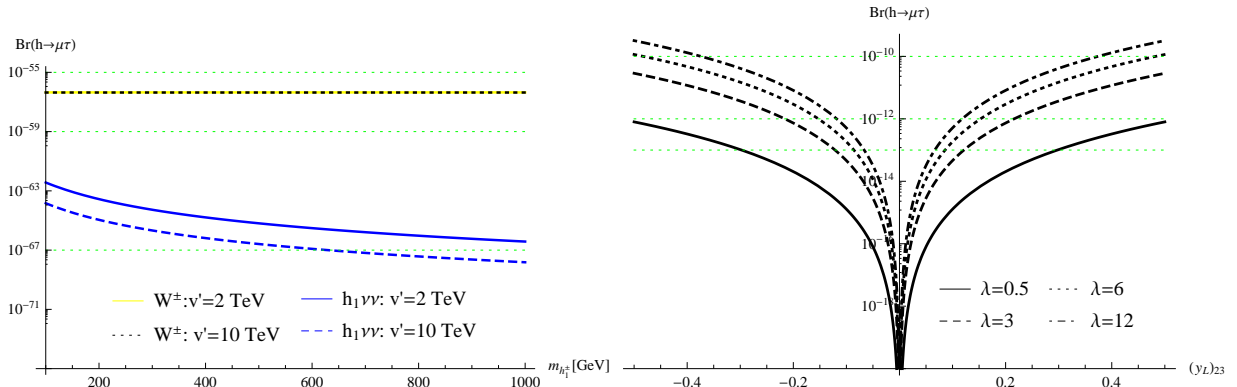


FIG. 2: Private contributions of only active neutrino loops to the  $\text{BR}(h \rightarrow \mu\tau)$ , where the left hand side is assumed that  $(y_L)_{23} = \frac{0.14 m_{h_1^\pm}}{1\text{TeV}}$  and the right hand side corresponds to  $v' = 10$  TeV.

points: the tiny contributions in the left panel and the significantly enhancement of the

contribution in the right panel. Similarly to many well-known models, the total contribution of electroweak loops, including  $W^\pm$  and their Goldstone bosons, is very suppressed because of the GIM mechanism, arising from the sum of two different flavors of external lepton fields:  $\sum_a U_{2a}^L U_{3a}^{L*} = 0$ . This sum cancels the largest terms of the contributions when they are expanded in terms of  $(m_{\nu_a}/m_W)^2$  series. Only terms containing factors  $(m_{\nu_a}/m_W)^2$  survive but they are very suppressed. In addition, this contribution does not depend on the  $v'$ , leading to the overlap lines in the left panel of the figure. The second contribution in the left panel comes from the  $\nu\nu h_1^\pm$  loops. Although the appearance of the Yukawa couplings  $(y_L)_{ij}$  removes the GIM mechanism, the contribution itself contains a factor of  $m_{\nu_a}^2$ ; therefore it is suppressed, too. It is even smaller than the electroweak-loop contribution because  $m_{h_1^\pm}$  is much larger than  $m_W$ . The  $\nu h_1^\pm h_1^\pm$  is much enhanced because of the presence of both large coupling  $\lambda_{hh_1^\pm h_1^\pm}$  and  $y_L$ . In the model considered,  $|(y_L)_{23}|$  is much constrained from (22), where the  $m_{h_1^\pm} = 3$  TeV gives the small  $|(y_L)_{23}| < 0.45$ . Also, the  $\lambda_{hh_1^\pm h_1^\pm}$ , as functions of  $v'$  ( $< 10$  TeV) and Higgs self-couplings ( $< 4\pi$ ), does not allow large values of  $\text{Br}(h \rightarrow \mu\tau)$ . Then the contribution of the  $\nu\nu h_1^\pm$  loop to the  $\text{Br}(h \rightarrow \mu\tau)$  is not larger than  $10^{-10}$ . In any case, this provides a hint for enhancing the contribution from active neutrino loops, for example in models with a four-loop (or higher) neutrino mass and small  $m_{h_1^\pm}$ , where the constraint of  $y_L$  may be released. This deserves further study.

The main contributions of exotic leptons to the  $\text{Br}(h \rightarrow \mu\tau)$  come from the two diagrams  $Nh_2^\pm h_2^\pm$  and  $h_2^\pm NN$ . They do not depend on  $y_L$  but depend strongly on  $|(y_R)_{ij}|$  with  $\{i, j\} = \{2, 3\}$ . As illustrated in Figs. 3 and 4, the two contributions have common properties. They increase with a decreasing  $m_{h_2^\pm}$ , in the same behaviour shown in the right panel of the Fig. 3. Most important is that they are enhanced strongly with an increasing  $|y'_R|$ , see the two left panels of the two figures.

On the other hand, these two contributions behave in opposing ways with the changes of  $v'$  and  $m_{N_2}(m_{N_3})$ . The contribution from  $m_{N_1}$  mediation is ignored because of a very small  $|(y_R)_{i1,1i}|$  ( $i = 1, 2, 3$ ). The  $Nh_2^\pm h_2^\pm$  mediation relates with the Higgs-self coupling; hence its contribution is large with a small  $m_{N_2}$  and a large  $v'$ . See the figure 3 with its four fixed values of  $v'$ , where the largest corresponds to  $v' = 10$  TeV. For the  $h_2^\pm NN$  loops, their analytic expression contains  $m_{N_a}^2/v'$  factors separate from the  $C$  functions. Therefore, the small  $v'$  and large  $m_{N_a}$  will give large contributions; see in the right panel of Fig. 4.

All of the above discussions suggest that the  $\text{Br}(h \rightarrow \mu\tau)$  will be large with small singly

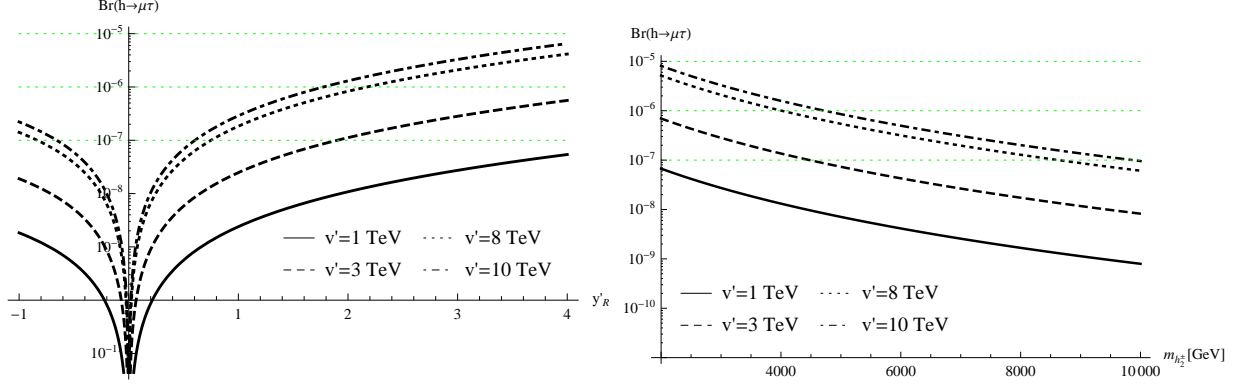


FIG. 3: Private contributions of only  $Nh_2^\pm h_2^\pm$  loops, i.e., diagrams 1 d), to the  $\text{Br}(h \rightarrow \mu\tau)$  as functions of  $y'_R$  ( $m_{h_2}$ ) in the left (right) panel.

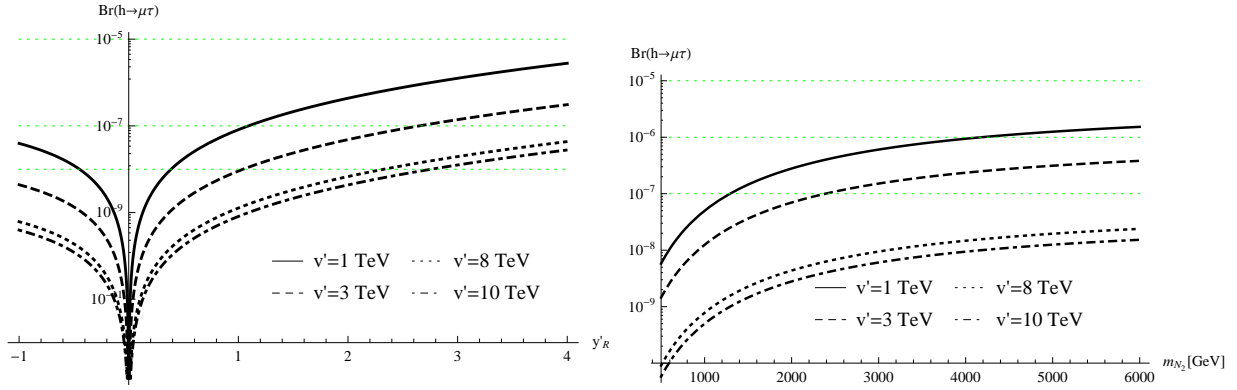


FIG. 4: Private contributions of only  $h_2^\pm NN$  loops, i.e., diagrams 1 f), to the  $\text{Br}(h \rightarrow \mu\tau)$  as functions of  $y'_R$  ( $m_{N_2}$ ) in the left (right) panel.

charged Higgs masses and large values of all of the following parameters: the coupling  $|y'_R|$ ,  $|(y_L)_{23}|$  and Higgs-self coupling  $\lambda$ . The dependence of the BR on  $v'$  and  $m_{N_{2,3}}$  is a bit complicated. The dependence on the mixing angle  $\alpha$  of two CP-even neutral Higgses should also be mentioned. Figure 5 shows more precisely the variations of the  $\text{Br}(h \rightarrow \mu\tau)$  on some particular free parameters. We can realize that  $y'_R$  affects the change of this BR the most significantly. It is larger than  $10^{-6}$  only when  $|y'_R| \geq 2$ .  $\text{Br}(h \rightarrow \mu\tau)$  does not depend on the signs of  $s_\alpha$  and  $y'_R$ , but it does depend significantly on the absolute values of these parameters. With  $\lambda = 4$ ,  $m_{h_2^\pm} = 3$  TeV and  $s_\alpha = 0.3$ , the  $\text{Br}(h \rightarrow \mu\tau)$  can reach  $10^{-5}$  when all of these conditions are satisfied:  $|y'_R| \geq 3$ ,  $v' \geq 8$  TeV and  $m_{N_{2,3}}$  are small enough.

Finally, if the doubly charged Higgs bosons are heavy enough, the investigation shown in [17] may allow the presence of light, singly charged Higgs bosons  $h_2^\pm$ , for instance with the

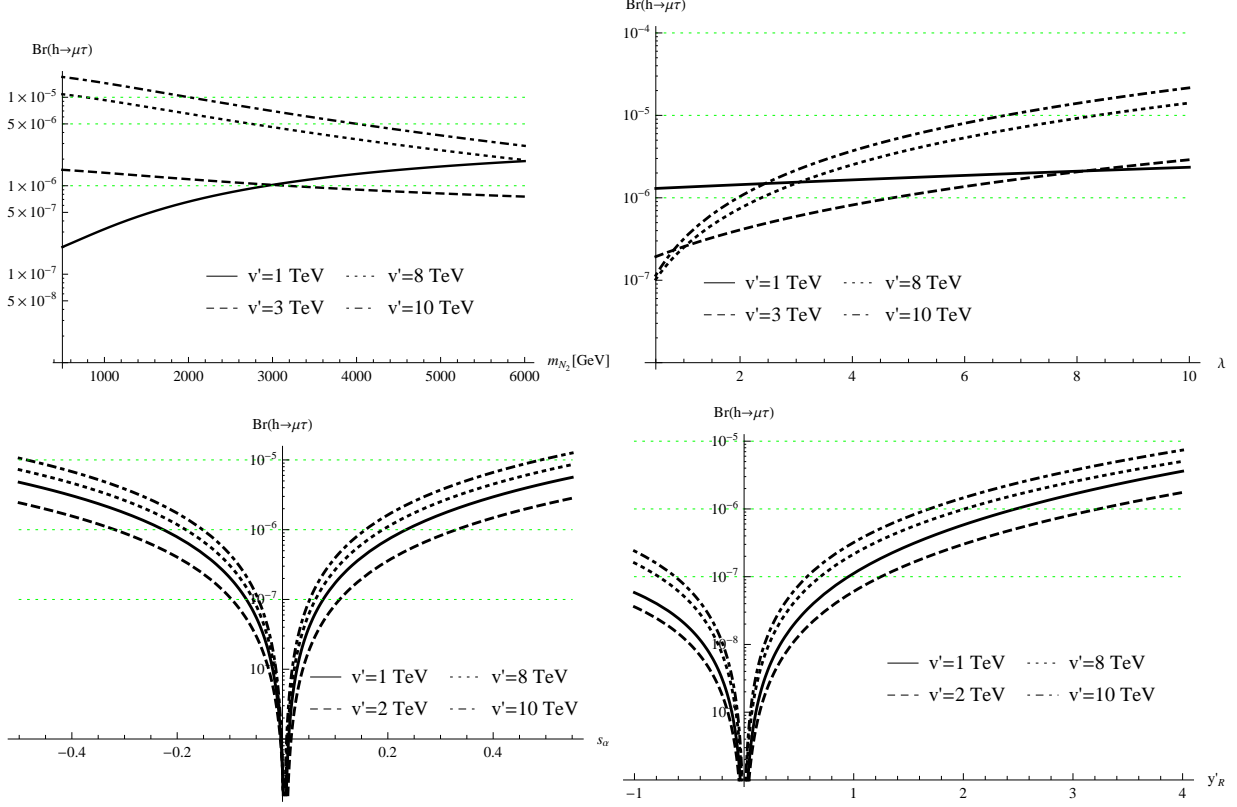


FIG. 5: The total contribution to the  $\text{Br}(h \rightarrow \mu\tau)$  as functions of single parameters:  $m_{N_2}$ ,  $\lambda$ ,  $s_\alpha$  or  $y'_R$ .

mass of 1 TeV. Besides, if we define  $m_{N_2} = m_{N_3}/2 = f \times v'$ , adopting also that  $s_\alpha = 0.3$  gives  $v' \leq 9$  TeV [17, 18]. In addition, if  $y'_R = 4$ , then  $\text{BR}(h \rightarrow \mu\tau)$  can reach the value of  $10^{-4}$ , very close to the value noted by CMS. This conclusion is illustrated in Fig. 6, where the  $f$  parameter is scanned in the range of  $0.1 \leq f \leq 6$  ( $0.1 \leq f \leq 2$ ) in the left (right) panel. From our numerical calculation, the values of  $y'_R = 4$  and  $m_{h_1^\pm} = 1$  TeV are the smallest ones to obtain values about  $10^{-4}$  of  $\text{Br}(h \rightarrow \mu\tau)$ . Accordingly, this large value lies in the only region where  $v'$  is larger than 4 TeV, while  $m_{N_2}$  and  $m_{N_3}$  should be as small as possible. In this calculation,  $m_{N_2} \geq 400$  GeV and the  $Nh_2^\pm h_2^\pm$  contribution is dominant. With small  $v' \simeq 1$  TeV, the value around  $10^{-5}$  of the  $\text{Br}(h \rightarrow \mu\tau)$  may occur if  $m_{N_{2,3}}$  are large enough:  $m_{N_2} \geq 3v'$ . And the  $\text{Br}(h \rightarrow \mu\tau)$  is unchanged, with general condition  $m_{N_2} \leq m_{N_3}$ , instead of the  $M_{N_2} = 1/2 M_{N_3}$  one used above. In this case, the large LFVHD corresponds to the dominant contribution from  $h_2^\pm NN$  diagram. The heavy neutrinos may be detected in the future lepton colliders [28].

The singly charged Higgses are now being searched by experiments [22, 25]. However,



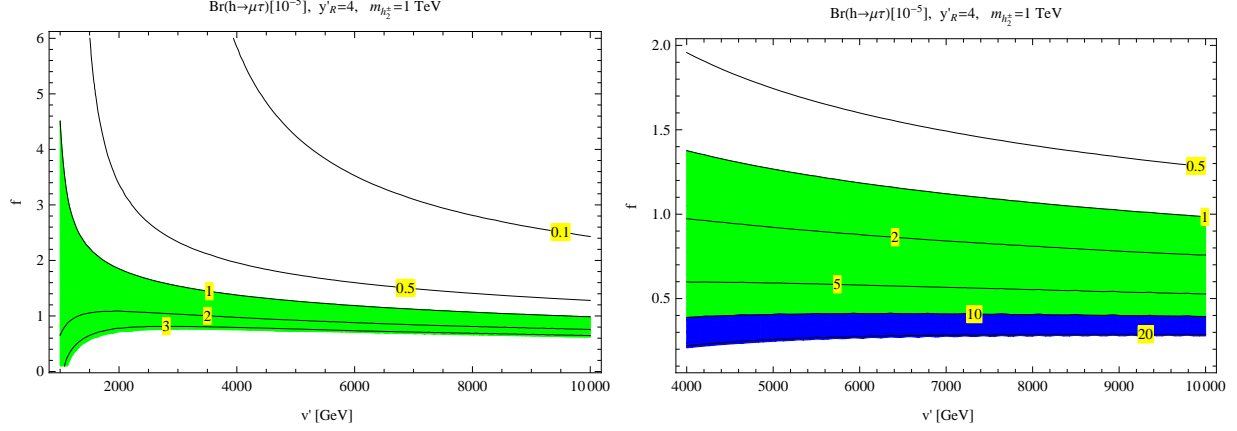


FIG. 6: Contour plots of  $\text{Br}(h \rightarrow \mu\tau)$  as functions of  $v'$  and  $f = m_{N_2}/v'$ . The green and blue regions respectively satisfy  $10^{-5} \leq \text{Br}(h \rightarrow \mu\tau) \leq 10^{-4}$  and  $\text{Br}(h \rightarrow \mu\tau) \geq 10^{-4}$ .

in the model considered, none of them are their targets, because they only couple with leptons. In addition, the  $h_2^\pm$  hold negative parities and hence cannot decay to only normal leptons. We also believe that the condition (22) is enough to guarantee for constraints of LFV processes without the need for singly charged Higgs bosons that are too heavy. Hence, the mass around 1 TeV of  $h_2^\pm$  is reasonable for stability of the lightest  $N_{R_1}$  as a DM candidate.

## V. CONCLUSION

Radiative neutrino mass models are interesting ones for explaining the tiny masses of active neutrinos. In the model introduced in [17], the new parameters that generate neutrino masses radiatively are strongly constrained from recent experimental data such as neutrino oscillations, rare decay of charged leptons, gauge coupling universalities and other constraint from LHC Higgs boson physics. Therefore these parameters are very predictive for other phenomenologies such as dark matter, LFVHD, etc... In this work, we have shown that in the allowed region with heavy, singly charged Higgs bosons, where their masses are around 3 TeV, the BR of the LFVHD can reach a value of  $10^{-5}$ . If the mass of the  $h_2^\pm$  is around 1 TeV, the  $\text{Br}(h \rightarrow \mu\tau)$  can reach  $10^{-4}$ . The additional necessary conditions are  $v'$ ,  $|s_\alpha|$  and all amplitudes of Yukawa and trilinear Higgs couplings must be large enough. For example, with largest  $s_\alpha = 0.3$  to satisfy the LHC Higgs boson constraint,  $v'$  should be in the range of 8 – 9 TeV. Also,  $\lambda_{\Sigma h_2} \geq 4$  and  $|y'_R| \geq 3(4)$  for largest values  $10^{-5}(10^{-4})$  of the  $\text{Br}(h \rightarrow \mu\tau)$ . The masses of the two heavy exotic neutrinos should be small, around 400 GeV for the

lighter  $m_{N_2}$ . In the model under consideration, a large  $\text{BR}(h \rightarrow \mu\tau)$  will lead to necessary consequences: the doubly charged Higgses  $k^{\pm\pm}$  must be heavy and the  $\text{BR}(h \rightarrow e\tau)$  must be very small. The latter can be explained from the constraint of  $\mu \rightarrow e\gamma$ . This requires a large  $y'_R \simeq \mathcal{O}(1)$  to give very small  $|(y_R)_{13,31}| \leq 10^{-3}$ , implying that  $\text{BR}(h \rightarrow e\tau) \leq 10^{-10}$ , much smaller than recent sensitivity of experiments [4]. One more interesting property is that the contribution of virtual active neutrinos may be large if the upper bound (22) is ignored. When this bound is considered the private contributions of  $\nu h_1^\pm h_1^\pm$  to  $\text{Br}(h \rightarrow \mu\tau)$  are around  $10^{-10}$ , which is much larger than values predicted by canonical seesaw models and [10]. In models with more than three-loop neutrino mass such as [27], the bound (22) may be released. We then guess that the active neutrino contributions can reach  $10^{-7}$  or higher, and hence they should not be ignored. Inclusion, more precise predictions can be worked out after the updating of data from the near future experiments.

### Acknowledgments

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### Appendix A: One loop Passarino-Veltman functions

Calculation in this section relates to one-loop diagrams in Fig. 1. The analytic expressions of the PV function are given in [10] and the needed formulas will be summarized in this appendix. We would like to stress that these PV functions were derived from the general form given in [26], using only the conditions of very small masses of tauons and muons. They are consistent with [21]. The denominators of the propagators are denoted as  $D_0 = k^2 - M_0^2 + i\delta$ ,  $D_1 = (k - p_1)^2 - M_1^2 + i\delta$  and  $D_2 = (k + p_2)^2 - M_2^2 + i\delta$ , where  $\delta$  is infinitesimally a positive real quantity. The scalar integrals are defined as

$$\begin{aligned} B_0^{(1)} &\equiv \frac{(2\pi\mu)^{4-D}}{i\pi^2} \int \frac{d^D k}{D_0 D_1}, & B_0^{(2)} &\equiv \frac{(2\pi\mu)^{4-D}}{i\pi^2} \int \frac{d^D k}{D_0 D_2}, \\ B_0^{(12)} &\equiv \frac{(2\pi\mu)^{4-D}}{i\pi^2} \int \frac{d^D k}{D_1 D_2}, & C_0 &\equiv C_0(M_0, M_1, M_2) = \frac{1}{i\pi^2} \int \frac{d^4 k}{D_0 D_1 D_2}, \end{aligned} \quad (\text{A1})$$

where  $i = 1, 2$ . In addition,  $D = 4 - 2\epsilon \leq 4$  is the dimension of the integral;  $M_0, M_1, M_2$  are masses of virtual particles in the loop. The momenta satisfy conditions:  $p_1^2 = m_1^2$ ,  $p_2^2 = m_2^2$  and  $(p_1 + p_2)^2 = m_h^2$ . In this work, with  $m_1$  and  $m_2$  are respective masses of the muon and the tauon, and  $m_h$  is the SM-like Higgs boson mass. The tensor integrals are

$$\begin{aligned} B^\mu(p_i; M_0, M_i) &= \frac{(2\pi\mu)^{4-D}}{i\pi^2} \int \frac{d^D k \times k^\mu}{D_0 D_i} \equiv B_1^{(i)} p_i^\mu, \\ B^\mu(p_1, p_2; M_1, M_i) &= \frac{(2\pi\mu)^{4-D}}{i\pi^2} \int \frac{d^D k \times k^\mu}{D_1 D_2} \equiv B_1^{(12)} p_1^\mu + B_2^{(12)} p_2^\mu, \\ C^\mu &= C^\mu(M_0, M_1, M_2) = \frac{1}{i\pi^2} \int \frac{d^4 k \times k^\mu}{D_0 D_1 D_2} \equiv C_1 p_1^\mu + C_2 p_2^\mu, \end{aligned} \quad (\text{A2})$$

where  $B_{0,1}^{(i)}$  and  $C_{0,1,2}$  are PV functions. It is well known that  $C_i$  is finite while the remains are divergent. We define  $\Delta_\epsilon \equiv \frac{1}{\epsilon} + \ln 4\pi - \gamma_E + \ln \frac{\mu^2}{m_h^2}$  where  $\gamma_E$  is the Euler constant. The divergent parts of the above scalar factors can be determined as

$$\text{Div}[B_0^{(i)}] = \text{Div}[B_0^{(12)}] = \Delta_\epsilon, \text{Div}[B_1^{(1)}] = -\text{Div}[B_1^{(2)}] = \frac{1}{2}\Delta_\epsilon. \quad (\text{A3})$$

The finite parts of the PV-functions such as B-functions depend on the scale of  $\mu$  parameter with the same coefficient of the divergent parts.

The analytic formulas of the above PV-functions are

$$B_{0,1}^{(i)} = \text{Div}[B_{0,1}^{(i)}] + b_{0,1}^{(i)}, \quad B_{0,1,2}^{(12)} = \text{Div}[B_{0,1,2}^{(12)}] + b_{0,1,2}^{(12)}. \quad (\text{A4})$$

The expression of  $b_0^{(12)}$  is

$$b_0^{(12)} = \ln \frac{m_h^2 - i\delta}{M_1^2 - i\delta} + 2 + \sum_{k=1}^2 x_k \ln \left( 1 - \frac{1}{x_k} \right), \quad (\text{A5})$$

where  $x_k$ , ( $k = 1, 2$ ) are solutions of the equation

$$x^2 - \left( \frac{m_h^2 - M_1^2 + M_2^2}{m_h^2} \right) x + \frac{M_2^2 - i\delta}{m_h^2} = 0. \quad (\text{A6})$$

The  $C_0$  function was given in [10] consistent with that discussed on [21], namely

$$C_0 = \frac{1}{m_h^2} [R_0(x_0, x_1) + R_0(x_0, x_2) - R_0(x_0, x_3)], \quad (\text{A7})$$

where

$$R_0(x_0, x_i) \equiv Li_2\left(\frac{x_0}{x_0 - x_i}\right) - Li_2\left(\frac{x_0 - 1}{x_0 - x_i}\right), \quad (\text{A8})$$

$Li_2(z)$  is the di-logarithm function,  $x_{1,2}$  are solutions of the equation (A6), and  $x_{0,3}$  are given as

$$x_0 = \frac{M_2^2 - M_0^2}{m_h^2}, \quad x_3 = \frac{-M_0^2 + i\delta}{M_1^2 - M_0^2}. \quad (\text{A9})$$

For simplicity of calculation we use approximate forms of PV functions where  $p_1^2, p_2^2 \rightarrow 0$ , namely,

$$\begin{aligned} b_0^{(i)} &= 1 - \ln \frac{M_i^2}{m_h^2} + \frac{M_0^2}{M_0^2 - M_i^2} \ln \frac{M_i^2}{M_0^2}, \\ b_1^{(1)} &= -\frac{1}{2} \ln \frac{M_1^2}{m_h^2} - \frac{M_0^4}{2(M_0^2 - M_1^2)^2} \ln \frac{M_0^2}{M_1^2} + \frac{(M_0^2 - M_1^2)(3M_0^2 - M_1^2)}{4(M_0^2 - M_1^2)^2}, \\ b_1^{(2)} &= \frac{1}{2} \ln \frac{M_2^2}{m_h^2} + \frac{M_0^4}{2(M_0^2 - M_2^2)^2} \ln \frac{M_0^2}{M_2^2} - \frac{(M_0^2 - M_2^2)(3M_0^2 - M_2^2)}{4(M_0^2 - M_2^2)^2}, \\ b_0^{(12)} &= \ln \frac{m_h^2 - i\delta}{M_1^2 - i\delta} + 2 + \sum_{k=1}^2 x_k \ln \left( 1 - \frac{1}{x_k} \right), \\ C_1 &= \frac{1}{m_h^2} \left[ b_0^{(1)} - b_0^{(12)} + (M_2^2 - M_0^2)C_0 \right], \quad C_2 = -\frac{1}{m_h^2} \left[ b_0^{(2)} - b_0^{(12)} + (M_1^2 - M_0^2)C_0 \right]. \end{aligned}$$

If  $M_1 = M_2$ , it can be seen that  $b_1^{(1)} = -b_1^{(2)}$ ,  $b_0^{(1)} = b_0^{(2)}$  and  $C_1 = -C_2$ .

## Appendix B: Form factors for LFVHD in t' Hooft-Feynman gauge

In this section we will list all factors for calculating LFVHD in the model considered. The calculation is done in the t' Hooft-Feynman gauge. The Feynman rules are given in Fig. 7. These factors were cross-checked using FORM [24].

Contribution from Fig. 1 a):

$$i\mathcal{M}_{(1a)} = \frac{ig^2 G_{hWW}}{16\pi^2} \sum_{a=1}^3 U_{2a}^L U_{3a}^{L*} \times ([\bar{u}_1 P_L v_2] \times E_L^{\nu WW} + [\bar{u}_1 P_R v_2] \times E_R^{\nu WW}), \quad (\text{B1})$$

where  $E_L^{\nu WW} = m_1 m_W C_1$ ,  $E_R^{\nu WW} = -m_2 m_W C_2$  and  $C_i \equiv C_i(m_{\nu_a}, m_W, m_W)$ .

Contribution from Fig. 1 b):

$$i\mathcal{M}_{(b)} = \frac{ig G_{hWG_w}}{16\pi^2} \sum_{a=1}^3 U_{2a}^L U_{3a}^{L*} \times ([\bar{u}_1 P_L v_2] E_L^{\nu WG_w} + [\bar{u}_1 P_R v_2] E_R^{\nu WG_w}), \quad (\text{B2})$$

where

$$\begin{aligned} E_L^{\nu WG_w} &= -\frac{m_1 m_2^2}{v} (2C_1 - C_2), \\ E_R^{\nu WG_w} &= -\frac{m_2}{v} \left( B_0^{(12)} + m_{\nu_a}^2 C_0 - m_1^2 C_1 + 2m_2^2 (-C_1 + C_2) + 2m_h^2 C_1 \right), \end{aligned} \quad (\text{B3})$$

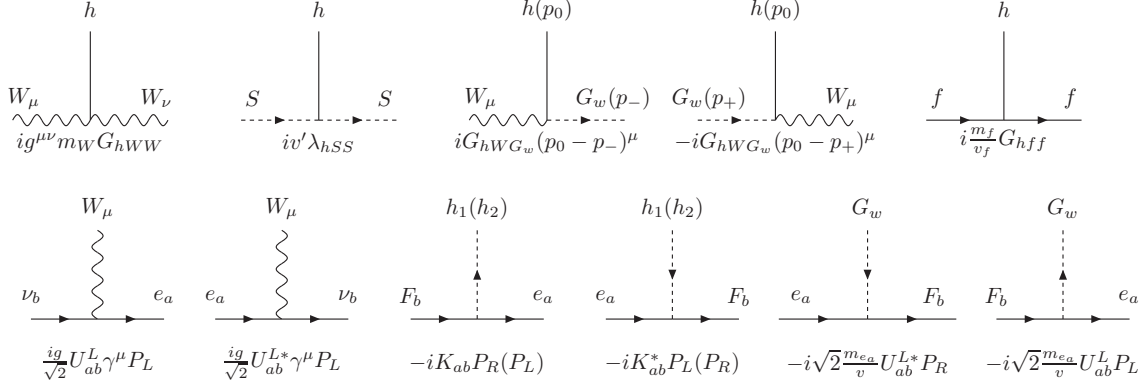


FIG. 7: Feynman rules for the  $h \rightarrow \mu^\pm \tau^\mp$  in the t' Hooft-Feynman gauge. Notations: i)  $S = G_w, h_1^\pm, h_2^\pm$ ; ii)  $K_{ab} = (y_L^T U^L)_{ab}$  and  $(y_R^T)_{ab}$  for active and exotic neutrinos, respectively; iii)  $f = e_a, \nu_a, N_a, F_a = \nu_a, N_a$ .

and  $B_0^{(12)} \equiv B_0^{(12)}(m_W, m_W)$ ,  $C_i \equiv C_i(m_{\nu_a}, m_W, m_W)$ .

Contribution from Fig. 1 c):

$$i\mathcal{M}_{(c)} = \frac{igG_{hWG_w}}{16\pi^2} \sum_{a=1}^3 U_{2a}^L U_{3a}^{L*} \times ([\bar{u}_1 P_L v_2] E_L^{\nu G_w W} + [\bar{u}_1 P_R v_2] E_R^{\nu G_w W}), \quad (\text{B4})$$

where

$$\begin{aligned} E_L^{\nu G_w W} &= -\frac{m_1}{v} \left( B_0^{(12)} + m_{\nu_a}^2 C_0 + 2m_1^2(-C_1 + C_2) + m_2^2 C_2 - 2m_h^2 C_2 \right), \\ E_R^{\nu G_w W} &= -\frac{m_1^2 m_2}{v} (C_1 - 2C_2), \end{aligned} \quad (\text{B5})$$

and  $B_0^{(12)} \equiv B_0^{(12)}(m_W, m_W)$ ,  $C_i \equiv C_i(m_{\nu_a}, m_W, m_W)$ .

Contribution from Fig. 1 d):

$$i\mathcal{M}_{(d)} = \frac{i(v' \lambda_{hSS})}{16\pi^2} \sum_{a=1}^3 V_{2a} V_{3a}^* \times ([\bar{u}_1 P_L v_2] E_L^{FSS} + [\bar{u}_1 P_R v_2] E_R^{FSS}), \quad (\text{B6})$$

where  $C_i \equiv C_i(m_a, m_S, m_S)$ ,  $V = \{U^L, K, y_R^T\}$ ,  $m_a = \{m_{\nu_a}, m_{N_a}\}$  for  $S = \{G_w, h_1, h_2\}$ ; and

$$\begin{aligned} E_L^{\nu G_w G_w} &= \frac{m_1 m_2^2}{v^2} 2C_2, & E_R^{\nu G_w G_w} &= \frac{m_1^2 m_2}{v^2} (-2C_1), \\ E_L^{\nu h_1 h_1} &= -m_1 C_1, & E_R^{\nu h_1 h_1} &= m_2 C_2; & E_L^{N h_2 h_2} &= m_2 C_2, & E_R^{N h_2 h_2} &= -m_1 C_1, \end{aligned} \quad (\text{B7})$$

Contribution from Fig. 1 e):

$$i\mathcal{M}_{(e)} = \frac{ig^2 G_{h\nu\nu}}{16\pi^2} \sum_{a=1}^3 U_{2a}^L U_{3a}^{L*} ([\bar{u}_1 P_L v_2] E_L^{W\nu\nu} + [\bar{u}_1 P_R v_2] E_R^{W\nu\nu}), \quad (\text{B8})$$

where

$$E_L^{W\nu\nu} = -\frac{m_1 m_{\nu_a}^2}{v'} (C_0 - 2C_1), \quad E_R^{W\nu\nu} = -\frac{m_2 m_{\nu_a}^2}{v'} (C_0 + 2C_2), \quad (\text{B9})$$

and  $C_i \equiv C_i(m_W, m_{\nu_a}, m_{\nu_a})$ .

Contribution from Fig. 1 f):

$$i\mathcal{M}_{(f)} = \frac{iG_{hFF}}{16\pi^2} \sum_{a=1}^3 V_{2a} V_{3a}^* ([\bar{u}_1 P_L v_2] E_L^{SFF} + [\bar{u}_1 P_R v_2] E_R^{SFF}), \quad (\text{B10})$$

where  $C_i \equiv C_i(m_S, m_a, m_a)$ ,  $V = \{U^L, K\}$ ,  $m_a = \{m_{\nu_a}, m_{N_a}\}$  for  $S = \{G_w, h_1, h_2\}$ ; and

$$\begin{aligned} E_L^{G_w\nu\nu} &= -\frac{m_1 m_{\nu_a}^2}{v'} \left[ \frac{m_2^2}{v^2} 2(C_0 + 2C_2) \right], & E_R^{G_w\nu\nu} &= -\frac{m_2 m_{\nu_a}^2}{v'} \left[ \frac{m_1^2}{v^2} 2(C_0 - 2C_1) \right], \\ E_L^{h_1\nu\nu} &= -\frac{m_1 m_{\nu_a}^2}{v'} (C_0 - 2C_1), & E_R^{h_1\nu\nu} &= -\frac{m_2 m_{\nu_a}^2}{v'} (C_0 + 2C_2), \\ E_L^{h_2 NN} &= -\frac{m_2 m_{N_a}^2}{v'} (C_0 + 2C_2), & E_R^{h_2 NN} &= -\frac{m_1 m_{N_a}^2}{v'} (C_0 - 2C_1), \end{aligned} \quad (\text{B11})$$

Contribution from sum of Figs. 1 g) and 1 h):

$$i\mathcal{M}_{(gh)} = \frac{ig^2 G_{hee}}{16\pi^2} \sum_{a=1}^3 U_{2a}^L U_{3a}^{L*} ([\bar{u}_1 P_L v_2] E_L^{\nu W} + [\bar{u}_1 P_R v_2] E_R^{\nu W}), \quad (\text{B12})$$

where

$$E_L^{\nu W}(m_F, m_W) = -\frac{m_1 m_2^2}{(m_1^2 - m_2^2)v} [B_1^{(1)} + B_1^{(2)}], \quad E_R^{\nu W}(m_F, m_W) = \frac{m_1}{m_2} E_L^{\nu W}, \quad (\text{B13})$$

$$B_1^{(i)} \equiv B_1^{(i)}(m_{\nu_a}, m_W).$$

Contribution from sum of Figs. 1 i) and 1 k):

$$i\mathcal{M}_{(ik)} = \frac{iG_{hee}}{16\pi^2} \sum_{a=1}^3 V_{2a} V_{3a}^* ([\bar{u}_1 P_L v_2] E_L^{FS} + [\bar{u}_1 P_R v_2] E_R^{FS}), \quad (\text{B14})$$

where  $B_1^{(i)} \equiv B_1^{(i)}(m_a, m_S)$ ,  $V = \{U^L, K, y_R^T\}$ ,  $m_a = \{m_{\nu_a}, m_{N_a}\}$  for  $S = \{G_w, h_1, h_2\}$ ; and

$$\begin{aligned} E_L^{\nu G_w} &= -\frac{2m_1^3 m_2^2}{(m_1^2 - m_2^2)v^3} (B_1^{(1)} + B_1^{(2)}), & E_R^{\nu G_w} &= \frac{m_2}{m_1} E_L^{\nu G_w}, \\ E_L^{\nu h_1} &= \frac{-m_1 m_2^2}{(m_1^2 - m_2^2)v} (B_1^{(1)} + B_1^{(2)}), & E_R^{\nu h_1} &= \frac{m_1}{m_2} E_L^{\nu h_1}, \\ E_L^{Nh_2} &= -\frac{m_1^2 m_2}{(m_1^2 - m_2^2)v} (B_1^{(1)} + B_1^{(2)}), & E_R^{Nh_2} &= \frac{m_2}{m_1} E_L^{Nh_2}, \end{aligned} \quad (\text{B15})$$

The divergence cancellation of the total amplitudes of the LFV decays is proved as follows.

The divergences appear only in the expressions listed in (B3), (B5), (B13) and (B15). The

expressions in (B3) and (B5) will vanish after inserting them into (B2) and (B4), where the GIM mechanism works. The divergences in each contribution given in (B13) and (B15) cancel each other because  $\text{Div}[B_1^{(1)}] = -\text{Div}[B_1^{(2)}]$ . Furthermore, the limit  $p_1^2, p_2^2 \rightarrow 0$  results in  $B_1^{(1)} = -B_1^{(2)}$ ; therefore all of the aforementioned contributions are very suppressed.

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- [1] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. **B 716**, 1 (2012), arXiv:1207.7214.
  - [2] V. Khachatryan *et al.* (CMS Collaboration), Phys. Lett. **B 716**, 30 (2012), arXiv:1207.7235.
  - [3] V. Khachatryan *et al.* (CMS Collaboration), Phys. Lett. **B 749**, 337 (2015); G. Aad *et al.* (ATLAS Collaboration), JHEP **1511** (2015) 211, arXiv: hep-ex/1508.03372.
  - [4] CMS Collaboration, (2015), CMS-PAS-HIG-14-040, <http://inspirehep.net/record/1388810/>
  - [5] A. Pilaftsis, Phys. Lett. **B 285**, 68 (1992); J. G. Körner, A. Pilaftsis, K. Schilcher, Phys. Rev. D **47** (1993) 1080; E. Arganda, A. M. Curiel, M. J. Herrero, D. Temes, Phys. Rev. D **71**, 035011 (2005), arxiv: hep-ph/0407302.
  - [6] E. Arganda, M. J. Herrero, X. Marcano and C. Weiland, Phys. Rev. D **91**, 015001 (2015).
  - [7] A. Brignole, A. Rossi, Phys. Lett. **B 66**, 217 (2003), arXiv:hep-ph/0304081; A. Brignole, A. Rossi, Nucl. Phys. **B 701**, 3 (2004), arXiv:hep-ph/0404211; M. Arana-Catania, E. Arganda, M. J. Herrero, JHEP **1309**, 160 (2013); JHEP **1510**, 192 (2015); E. Arganda, M. J. Herrero, X. Marcano, C. Weiland, Phys. Rev. D **93**, 055010 (2016) , arXiv: hep-ph/1508.04623; P. T. Giang, L. T. Hue, D. T. Huong and H. N. Long, Nucl. Phys. **B 864** (2012) 85, [arXiv:1204.2902(hep-ph)]; L. T. Hue, D. T. Huong, H. N. Long, H. T. Hung, N. H. Thao, Prog. Theor. Exp. Phys. **113B05** (2015), arXiv:1404.5038 [hep-ph]; D. T. Binh, L. T. Hue, D. T. Huong, H. N. Long, Eur. Phys. J. **C 74** (2014) 2851, [arXiv:1308.3085(hep-ph)]; E. Arganda, M. J. Herrero, R. Morales and A. Szynekman, JHEP **1603**, 055 (2016), arxiv: hep-ph/1510.04685; J. L. Diaz-Cruz, JHEP **0305**, 036 (2003) arXiv: hep-ph/0207030].
  - [8] A. Crivellin, G. D'Ambrosio, J. Heeck, Phys. Rev. Lett. **114** (2015) 151801; N. Bizot, S. Davidson, M. Frigerio, J. L. Kneur, JHEP **1603** (2016) 073.
  - [9] N. Bizot, S. Davidson, M. Frigerio, J. -L. Kneur, JHEP **1603** (2016) 073; F. J. Botella, G. C. Branco, M. Nebot, M. N. Rebelo, Eur. Phys. J. **C 76** (2016), 161; S. Kanemura, T. Ota, T. Shindou and K. Tsumura, Phys. Rev. D **73**, 016006 (2006), hep-ph/0505191; M. Arroyo, J. L. Diaz-Cruz, E. Diaz and J. A. Orduz-Ducuara, arXiv: hep-ph/1306.2343; D. Das and A.

- Kundu, Phys. Rev. D **92**, 015009 (2015), arXiv: hep-ph/1504.01125; X. F Han, L. Wang, J. M. Yang, Phys. Lett. **B 757** (2016) 537.
- [10] L. T. Hue, H. N. Long, T. T. Thuc and T. Phong Nguyen, Nucl. Phys. **B 907** (2016) 37, arXiv: hep-ph/1512.03266.
- [11] S. Baek, Z.-F. Kang, JHEP **1603** (2016) 106; S. Baek, K. Nishiwaki, Phys. Rev. D **93** (2016), 015002.
- [12] K. Cheung, W. Y. Keung, P. Y. Tseng, Phys. Rev. D **93** (2016), 015010.
- [13] W. Altmannshofer, S. Gori, A. L. Kagan, L. Silvestrini, J. Zupan, Phys. Rev. D **93** (2016), 031301; X. G. He, J. Tandean, Y. J. Zheng, JHEP **1509** (2015) 093 ; I. Doršner, S. Fajfer, A. Greljo, J. F. Kamenik, N. Košnik, Ivan Nišandžić, JHEP **1506** (2015) 108; R. Harnik, J. Kopp and J. Zupan, JHEP **1303**, 026 (2013), arXiv: hep-ph/1209.1397; A. Celis, V. Cirigliano and E. Passemar, Phys. Rev. D **89**, 013008 (2014), arXiv: hep-ph/1309.3564; A. Dery, A. Efrati, Y. Nir, Y. Soreq and V. Susi, Phys. Rev. D **90**, 115022 (2014), arXiv: hep-ph/1408.1371; J. Heeck, M. Holthausen, W. Rodejohann and Y. Shimizu, Nucl. Phys. **B 896**, 281 (2015), arXiv: hep-ph/1412.3671; A. Crivellin, G. D'Ambrosio and J. Heeck, Phys. Rev. D **91**, 075006 (2015), arXiv: hep-ph/1503.03477; J. L. Diaz-Cruz and J. J. Toscano, Phys. Rev. D **62**, 116005 (2000), arXiv: hep-ph/9910233; A. Pilaftsis, Z. Phys. **C 55**, 275 (1992) ; L. D. Lima, C. S. Machado, R. D. Matheus, L. A. F. D. Prado, JHEP **1511**, 074 (2015), arXiv: hep-ph/1501.06923; I. d. M. Varzielas, O. Fischer, V. Maurer, JHEP **1508**, 080 (2015); C. F. Chang, C. H. V. Chang, C. S. Nugroho, T. C. Yuan, " *Lepton Flavor Violating Decays of Neutral Higgses in Extended Mirror Fermion Model* ", e-Print: arXiv:1602.00680; C. H. Chen, T. Nomura, " *Bound on LFV Higgs decays in a vectorlike lepton model and search for doubly charged lepton at the LHC* ", arXiv:1602.07519 [hep-ph]; K. Huitu, V. Keus, N. Koivunen, O. Lebedev, JHEP **1605** (2016) 026, arXiv:1603.06614.
- [14] M. Sher, K. Thrasher, Phys. Rev. D **93**, 055021 (2016).
- [15] S. Kanemura, K. Matsuda, T. Ota, T. Shindou, E. Takasugi and K. Tsumura, Phys. Lett. **B 599**, 83 (2004), hep-ph/0406316; G. Blankenburg, J. Ellis and G. Isidori, Phys. Lett. **B 712**, 386 (2012), arXiv:hep-ph/1202.5704; S. Davidson and P. Verrier, Phys. Rev. D **86**, 111701 (2012), arXiv: hep-ph/1211.1248; S. Bressler, A. Dery and A. Efrati, Phys. Rev. D **90**, 015025 (2014), arXiv: hep-ph/1405.4545; D. Aristizabal Sierra and A. Vicente, Phys. Rev. D **90**, 115004 (2014), arXiv: hep-ph/1409.7690; C. X. Yue, C. Pang and Y. C. Guo, J. Phys.



- G 42**, 075003 (2015), arXiv: hep-ph/1505.02209; S. Banerjee, B. Bhattacharjee, M. Mitra, M. Spannowsky, "*The Lepton Flavour Violating Higgs Decays at the HL-LHC and the ILC*", arXiv:1603.05952 [hep-ph]; I. Chakraborty, A. Datta, A. Kundu, "*Lepton flavor violating Higgs boson decay  $h \rightarrow \mu\tau$  at the ILC*", arXiv:1603.06681 [hep-ph].
- [16] Y. Fukuda, et al., Phys. Rev. Lett. **81**, 1562 (1998); S. Fukuda et al., Phys. Rev. Lett. **85** (2000) 3999.
- [17] K. Nishiwaki, H. Okada, and Y. Orikasa, Phys. Rev. D **92**, 093013 (2015).
- [18] S. Kanemura, K. Nishiwaki, H. Okada, Y. Orikasa, S. C. Park, R. Watanabe, "LHC 750 GeV Diphoton excess in a radiative seesaw model", arXiv:1512.09048 [hep-ph].
- [19] M. C. Gonzalez-Garcia, M. Maltoni, T. Schwetz, JHEP **1411** (2014) 052 .
- [20] H. K. Dreiner, H. E. Haber, S. P. Martin, Phys. Rept. **494**, 1 (2010).
- [21] D. Y. Bardin, G. Passarino, "*The Standard Model in the making: Precision study of the electroweak interactions*", Clarendon Press-Oxford, 1999.
- [22] K. A. Olive et al. (Particle Data Group), Chin. Phys. C **38**, 090001 (2014).
- [23] J. Herrero-Garcia, M. Nebot, N. Rius, and A. Sntamaria, Nulc. Phys. B **885**, 542 (2014).
- [24] J. A. M. Vermaseren, "*New features of FORM*", arxiv: math-ph/0010025.
- [25] ATLAS Collaboration (G. Aad *et al.*), JHEP **1206**, 039 (2012); ATLAS Collaboration (G. Aad *et al.*), JHEP **1603** (2016) 127.
- [26] G. 't Hooft and M. Veltman, Nucl. Phys. **B 153**, 365 (1979).
- [27] T. Nomura, H. Okada, Phys. Lett. **B 755** (2016) 306; T. Nomura, H. Okada, "*Four-loop Radiative Seesaw Model with 750 GeV Diphoton Resonance*", arXiv:1601.04516 [hep-ph].
- [28] S. Antusch, E. Cazzto, O. Fischer, JHEP **1604**, 189 (2016); S. Antusch, O. Fischer, JHEP **1505** (2015) 053; JHEP **1410**, 094 (2014).
- [29] Y. Mambrini, S. Profumo, and F. S. Queiroz, "*Dark Matter and Global Symmetries*", arXiv:1508.06635v1 [hep-ph].